

Light dark matter in the NMSSM: upper bounds on direct detection cross sections

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Abstract

In the Next-to-Minimal Supersymmetric Standard Model, a bino-like LSP can be as light as a few GeV and satisfy WMAP constraints on the dark matter relic density in the presence of a light CP-odd Higgs scalar. We study upper bounds on the direct detection cross sections for such a light LSP in the mass range $2 - 20$ GeV in the NMSSM, respecting all constraints from B-physics and LEP. The OPAL constraints on $e^+e^- \rightarrow \chi_1^0 \chi_i^0$ ($i > 1$) play an important rôle and are discussed in some detail. The resulting upper bounds on the spin-independent and spin-dependent nucleon cross sections are $\sim 10^{-42} \text{ cm}^{-2}$ and $\sim 4 \times 10^{-40} \text{ cm}^{-2}$, respectively. Hence the upper bound on the spin-independent cross section is below the DAMA and CoGeNT regions, but could be compatible with the two events observed by CDMS-II.

1 Introduction

The DAMA [1] and CoGeNT [2] dark matter detection experiments have reported events in excess of the expected background, which would be compatible with a WIMP mass of a few GeV. Also the CDMS-II experiment [3] has reported two events, which could be explained by a WIMP mass of $\gtrsim 10$ GeV (or background). On the other hand, exclusion limits from the Xenon10 [4], Xenon100 [5] and CDMS-Si [6] experiments set upper bounds on the spin independent detection cross sections for this mass range of a WIMP, which seem incompatible with the reported hints for a signal.

In any case, it is important to know whether specific models for dark matter with a WIMP mass of a few GeV can produce direct detection cross sections compatible with the reported excesses, and/or whether regions in the parameter space of such models can be tested by present and future exclusion limits.

Supersymmetric (Susy) extensions of the Standard Model are popular, amongst others, since they predict naturally (for unbroken R-parity, and for a neutral Lightest Supersymmetric Particle (LSP)) a candidate for dark matter, with a relic density compatible with WMAP constraints [7]. Within the Minimal Supersymmetric extension of the Standard Model (MSSM) four neutral fermions (neutralinos χ_i^0 , $i = 1 \dots 4$) exist, which are composed of the bino (superpartner of the $U(1)_Y$ gauge boson), the wino (superpartner of the W_μ^3 gauge boson) and two higgsinos (superpartners of neutral Higgs bosons). These states mix, and the lightest neutralino χ_1^0 , which is the lightest eigenvalue of the 4×4 mass matrix, will be the LSP (leaving aside the possibility of a sneutrino LSP).

Often the LSP is dominantly bino-like, whose mass $m_{\chi_1^0}$ is approximately given by the soft Susy breaking gaugino mass $\sim M_1$. Assuming unification of the three gaugino masses for the bino (M_1), the winos (M_2) and the gluino (M_3) at the scale of Grand Unification, M_1 is naturally the smallest among these mass terms at the electroweak scale. However, due to the lower bound of ~ 100 GeV on M_2 from the lower bound on chargino masses, the assumption of unification of the three gaugino masses implies $M_1 \gtrsim 50$ GeV and a similar lower bound on the mass of the LSP.

The assumption of unification of the three gaugino masses can be dropped, however, in that case M_1 and hence the LSP mass $m_{\chi_1^0}$ can be arbitrarily small. Then, on the other hand it can become difficult to satisfy the WMAP constraint on the dark matter relic density i.e., to reduce the dark matter relic density after the Big Bang to an acceptable value compatible with this constraint. To this end, dark matter annihilation processes have to be sufficiently effective. For a LSP mass $\gtrsim 50$ GeV the following pair annihilation processes can be relevant: exchange of Susy partners of fermions (sfermions, in particular sleptons) in the t-channel, and Z-exchange or Higgs-exchange in the s-channel (if the LSP has a sufficiently large higgsino component). In addition, neutralinos can co-annihilate with other sparticles if they have similar masses, but co-annihilation processes will not be relevant for a light LSP as considered here. In the MSSM, sufficiently effective dark matter annihilation processes impose constraints on a light LSP:

Considering LSP annihilation via slepton exchange in the t-channel, a lower bound $m_{\chi_1^0} \gtrsim 18$ GeV was derived in [8,9] from the lower bound of ~ 100 GeV on the slepton masses. (Relaxing this bound to $\gtrsim 80$ GeV for stau masses, one obtains $m_{\chi_1^0} \gtrsim 13$ GeV [8–10], unless the LSP mass is very small corresponding to hot dark matter [10].) Allow-

ing for LSP annihilation via CP-odd Higgs (A) exchange in the s-channel, a lower limit $m_{\chi_1^0} \gtrsim 6$ GeV was given in [11–14] from $m_A \gtrsim 90$ GeV for large values of $\tan\beta \gtrsim 25$. However, as noted in [15], this region of the parameter space of the MSSM is now strongly constrained by the bounds on $B_s \rightarrow \mu^+\mu^-$. A LSP with a mass in the 5 – 15 GeV range in the MSSM has been considered in [16] without, however, asking for a correct relic density.

In the Next-to-Minimal Supersymmetric Standard Model (NMSSM, see [17,18] for recent reviews), which can solve the μ -problem of the MSSM, the Higgs and neutralino sectors are extended by gauge singlet states. As noticed in [19–22], the mass of the LSP can be considerably smaller in the NMSSM than in the MSSM and can still be compatible with the WMAP constraint on the relic density. This is a consequence of a light CP-odd Higgs boson in the spectrum (on top of the CP-odd Higgs boson of the MSSM), which can be mostly singlet-like and which is not ruled out by LEP-constraints. Then, sufficiently large LSP annihilation cross sections via the exchange of this additional CP-odd Higgs boson in the s-channel may be possible even for a light LSP with mass of a few GeV.

A light LSP in the NMSSM could be a (dominantly) singlet-like state; in this case, however, its direct detection cross sections would be tiny. On the other hand, as in the MSSM, a light LSP in the NMSSM can originate from a small value of M_1 in which case it will be dominantly bino-like and can have larger direct detection cross sections. These have been estimated in [21], where also constraints on the corresponding parameter space from B-physics, LEP and Υ -physics were discussed. However, the corresponding points in the parameter space given as examples in [21] suffer from a negative effective μ -parameter (which is in conflict with the measured anomalous magnetic moment of the muon), and not all experimental constraints considered below are taken into account.

In view of the interest in a light LSP with a mass in the 2 – 20 GeV range, we find it appropriate to study upper bounds on its direct detection cross sections in the NMSSM. Direct detection cross sections in the NMSSM including WMAP constraints have been studied before in [23–29], but not for the LSP mass range considered here. Apart from WMAP constraints, we take care of a lengthy list of experimental constraints from B-physics (important for large $\tan\beta$ and/or relatively light charged and CP-odd Higgs bosons as relevant here), Υ -physics and LEP-constraints on neutralino production. Among the latter, OPAL limits on $e^+e^- \rightarrow \chi_1^0\chi_i^0$ ($i > 1$) turn out to be very important. Since these are also relevant for the MSSM, but have hardly been discussed before (a notable exception is [10]), we study their consequences in some detail. For the numerical analysis we use the code NMSSMTools [30,31] coupled to micrOMEGAs [32–34]. As a result we obtain upper bounds on the spin-independent and spin-dependent LSP-nucleon cross sections of $\sigma^{SI} \lesssim 10^{-42}$ cm² and $\sigma^{SD} \lesssim 4 \times 10^{-40}$ cm², varying somewhat with the LSP mass in the 2 – 20 GeV range. The maximal value for σ^{SI} is indeed near the estimate given in [21].

In the next section (2) we present the relevant parameters of the NMSSM and their impact on the LSP cross sections. In section (3), we discuss the relevant experimental constraints. The consequences of the OPAL constraints on $e^+e^- \rightarrow \chi_1^0\chi_i^0$ ($i > 1$) on the parameter space (implying a lower bound on μ_{eff}) are estimated in an analytic approximation, which reproduces well the full numerical results. Section (4) is devoted to our results and conclusions.

2 The NMSSM and the impact of its parameters on the LSP cross sections

In the NMSSM the μ parameter of the MSSM is replaced by a Yukawa coupling λ to a gauge singlet (super-) field S . Then, the vacuum expectation value (vev) s of the real scalar component of S generates an effective μ -term

$$\mu_{eff} = \lambda s . \quad (1)$$

Most of the time one studies the NMSSM with a scale invariant superpotential W which contains, apart from the Yukawa coupling of S to the MSSM-like Higgs doublet fields H_u and H_d , a trilinear term $\sim \frac{\kappa}{3}S^3$. Hence the Higgs mass term $\mu H_u H_d$ in W_{MSSM} is replaced by

$$W_{NMSSM} = \lambda S H_u H_d + \frac{\kappa}{3}S^3 + \dots . \quad (2)$$

(Occasionally one considers the so-called nMSSM [22, 25, 28] without the trilinear coupling $\sim \frac{\kappa}{3}S^3$, which is replaced by a tadpole-term $\sim \xi_F S$.) Compared to the MSSM, the gauge singlet superfield S adds additional degrees of freedom to the CP-even and CP-odd Higgs sectors as well as to the neutralino sector. Hence the spectrum contains

- 3 CP-even neutral Higgs bosons H_i , $i = 1, 2, 3$, which mix in general;
- 2 CP-odd neutral Higgs bosons A_1 and A_2 ;
- one charged Higgs boson H^\pm ;
- five neutralinos χ_i^0 , $i = 1 \dots 5$, which are mixtures of the bino, the neutral wino, the neutral higgsinos and the singlino;
- two charginos which are mixtures of the charged winos and the charged higgsinos.

Apart from the Susy generalisations of the Standard-Model-like gauge and Yukawa couplings and the superpotential in Eq. (2), the Lagrangian of the NMSSM contains soft Susy breaking terms in the form of gaugino masses M_1 , M_2 and M_3 for the bino, the winos and the gluino, respectively, mass terms for all scalars (squarks, sleptons, Higgs bosons including the singlet S) as well as trilinear scalar self-couplings as $\lambda A_\lambda S H_u H_d + \frac{\kappa}{3}A_\kappa S^3$, which reflect the trilinear couplings among the superfields in the superpotential.

Expressions for the mass matrices for all Higgs- and neutralino states can be found in [17, 18]; below we confine ourselves to those which are of relevance subsequently. Dropping the Goldstone mode, the 2×2 mass matrix for the CP-odd Higgs bosons \mathcal{M}_P^2 in the basis (A_{MSSM}, S_I) has the elements

$$\begin{aligned} \mathcal{M}_{P,11}^2 &= \frac{2\mu_{eff}(A_\lambda + \kappa s)}{\sin 2\beta} \equiv M_A^2 , \\ \mathcal{M}_{P,22}^2 &= \lambda(A_\lambda/s + 4\kappa)v_u v_d - 3\kappa A_\kappa s , \\ \mathcal{M}_{P,12}^2 &= \lambda(A_\lambda - 2\kappa s)v \end{aligned} \quad (3)$$

where v_u , v_d denote the vevs of H_u , H_d , respectively, $v = \sqrt{v_u^2 + v_d^2} \sim 174$ GeV and, as usual, $\tan \beta = v_u/v_d$. The matrix element $\mathcal{M}_{P,11}^2$ would be the mass squared of the MSSM-like CP-odd scalar A_{MSSM} , if the singlet sector were absent; subsequently we will denote it simply by M_A^2 . (This parameter can replace the parameter A_λ .) For any (possibly large) value of M_A^2 , \mathcal{M}_P^2 can have another small eigenvalue corresponding to an additional light CP-odd Higgs boson A_1 which is mostly singlet-like. This state will be relevant for the LSP annihilation cross section below.

The mass of the charged Higgs scalar is given by

$$M_{H^\pm}^2 = M_A^2 + v^2\left(\frac{g_2^2}{2} - \lambda^2\right); \quad (4)$$

note that it decreases with increasing λ . As is well known, too small values of M_{H^\pm} can cause disagreements between measurements and corresponding contributions to B-physics-observables as $b \rightarrow s\gamma$; this will be of importance below.

Notably for large M_A , one of the 3 CP-even Higgs bosons will have a mass close to M_A . In the MSSM, the corresponding CP-even state is denoted by H , and we will maintain this denomination. The spin-independent LSP-nucleon cross section will be dominated by the exchange of this CP-even scalar H , since its couplings to down-type quarks (particularly the strange quark) are enhanced for large values of $\tan \beta$.

Also, the mass of the charged Higgs scalar is close to M_A for large M_A ; then the states H , A_{MSSM} and H^\pm form a nearly degenerate SU(2) doublet. In fact this approximate degeneracy holds down to fairly low values of $M_A \sim 300$ GeV.

In the neutralino sector, the bino λ_1 and the neutral wino λ_2^3 mix with the neutral higgsinos ψ_d^0, ψ_u^0 and the singlino ψ_S , and generate a symmetric 5×5 mass matrix \mathcal{M}_0 . In the basis $\psi^0 = (-i\lambda_1, -i\lambda_2^3, \psi_d^0, \psi_u^0, \psi_S)$, \mathcal{M}_0 reads

$$\mathcal{M}_0 = \begin{pmatrix} M_1 & 0 & -\frac{g_1 v_d}{\sqrt{2}} & \frac{g_1 v_u}{\sqrt{2}} & 0 \\ & M_2 & \frac{g_2 v_d}{\sqrt{2}} & -\frac{g_2 v_u}{\sqrt{2}} & 0 \\ & & 0 & -\mu_{\text{eff}} & -\lambda v_u \\ & & & 0 & -\lambda v_d \\ & & & & 2\kappa_S \end{pmatrix}. \quad (5)$$

It can be diagonalized by an orthogonal real matrix N_{ij} such that the physical masses $m_{\chi_i^0}$ ordered in $|m_{\chi_i^0}|$ are real (but not necessarily positive). Denoting the 5 eigenstates by χ_i^0 , we have

$$\chi_i^0 = N_{ij} \psi_j^0. \quad (6)$$

Finally, the chargino masses are described by a 2×2 mass matrix containing M_2 and μ_{eff} as diagonal entries. The lower bound of ~ 103 GeV on the lightest chargino implies at least the constraint

$$\text{Min}\{M_2, |\mu_{\text{eff}}|\} \gtrsim 100 \text{ GeV} \quad (7)$$

(one can choose $M_2 > 0$ by convention).

Next, we discuss the dominant contribution to the spin-independent LSP-nucleon cross section σ^{SI} . Leaving aside scenarios with light squark masses of ~ 100 GeV (which are

difficult to reconcile with Tevatron constraints), σ^{SI} is dominated by the exchange of CP-even Higgs bosons, which couple mostly to the strange quark sea. Among the CP-even Higgs bosons, the coupling of the state H to down-type quarks (as the strange quark) increases with $\tan\beta$. Hence, although its mass is generally larger than the mass of the Standard-Model-like Higgs boson h , H -exchange provides the leading contribution to σ^{SI} for large values of $\tan\beta$. Then, the dominant component of H is given by H_d .

The dominant coupling of $H \sim H_d$ to the LSP is induced by the bino-higgsino-Higgs vertex $\sim g_1$ and hence proportional to $g_1 N_{11} N_{13}$, where N_{11} denotes the bino- and N_{13} the ψ_d^0 -higgsino-component of the LSP. All in all one finds

$$\sigma^{SI} \sim N_{11}^2 N_{13}^2 \frac{\tan^2 \beta}{m_H^4}, \quad (8)$$

which shows that the largest values of σ^{SI} are obtained for a large product $N_{11} N_{13}$, large $\tan\beta$ and low values of m_H .

The dominant contribution to the spin-dependent LSP-nucleon cross section σ^{SD} originates, as in the MSSM, from Z -exchange. At first sight one could imagine that, for a light CP-odd Higgs boson A_1 , its exchange could also give important contributions to σ^{SD} . However, a light A_1 is dominantly singlet-like and, moreover, the coupling of its doublet component to strange quarks is always tiny compared to the Z -boson coupling.

The coupling of the Z -boson to the LSP originates from the gauge couplings of the higgsino components ψ_u^0 and ψ_d^0 . Since no additional free parameters intervene, the spin-dependent cross section σ^{SD} is proportional to

$$\sigma^{SD} \sim (N_{13}^2 - N_{14}^2)^2. \quad (9)$$

Finally the LSP annihilation cross section σ_{ann} is dominated, for the LSP mass range 2 – 20 GeV under consideration, by the exchange of a light A_1 in the s-channel. The dominant contribution to the $A_1 \chi_1^0 \chi_1^0$ coupling is induced by the doublet component of A_1 and the bino-higgsino components of χ_1^0 as in the case of the $H \chi_1^0 \chi_1^0$ coupling above; the singlet components of A_1 and χ_1^0 play a minor rôle here. In any case one has (neglecting the finite width of A_1 and the velocity of χ_1^0 near the freeze-out temperature)

$$\sigma_{ann} \sim \frac{1}{(m_{A_1}^2 - 4m_{\chi_1^0}^2)^2}, \quad (10)$$

and hence σ_{ann} can be sufficiently large for suitable values of $m_{A_1}^2$, the lightest eigenvalue of \mathcal{M}_P^2 in Eq. (3).

3 Experimental constraints on the parameter space

In this section we discuss various constraints on the parameters of the NMSSM, notably (but not exclusively) from LEP and B-physics, separately in various subsections.

3.1 Constraints from sparticle and Higgs searches

As we have seen in Eq. (8), a large spin-independent detection cross section σ^{SI} requires bino-components N_{11} and higgsino-components N_{13} of the LSP. For small M_1 such that $m_{\chi_1^0}$ is in the 2 – 20 GeV range, the bino component of χ_1^0 is automatically large. However, a large higgsino component of χ_1^0 require relatively small values for μ_{eff} (below ~ 160 GeV) in the mass matrix \mathcal{M}_0 in Eq. (5). Consequently the neutralino states χ_2^0 and χ_3^0 (for $M_2, 2\kappa s > \mu_{eff}$) are higgsino-like with masses of the order of μ_{eff} . Then, the production process $e^+e^- \rightarrow \chi_1^0 \chi_i^0$ ($i = 2, 3$) was kinematically possible at LEP2, and corresponding limits from DELPHI [35] and OPAL [36] have to be taken into account.

The strongest limits come from OPAL at 208 GeV, where we can assume 100% Z^* branching ratios for the χ_i^0 decays (see Fig. 10 in [36]). Upper bounds on the cross section are given in 5 GeV-wide bins of $m_{\chi_i^0}$. Since we will find $m_{\chi_3^0} - m_{\chi_2^0} \sim 40$ GeV, the bounds apply for χ_2^0 and χ_3^0 separately. For $m_{\chi_1^0} < 20$ GeV, at least one of the χ_2^0 or χ_3^0 production cross sections (in association with χ_1^0) is bounded from above by 0.05 pb.

In principle, both Z^* -exchange in the s-channel and selectron exchange in the t-channel contribute to this cross section. However, the interference between these channels is positive, hence the most conservative bounds on the parameters are obtained by assuming heavy selectrons and that $e^+e^- \rightarrow \chi_1^0 \chi_i^0$ originates from Z^* -exchange only. The expression for $\sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0)$ is given, e.g., in [37] and can be written as

$$\begin{aligned} \sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0) &= \frac{(g_1^2 + g_2^2)^2}{32\pi(s - M_Z^2)^2} (N_{13}N_{i3} - N_{14}N_{i4})^2 \left(\frac{1}{4} - \sin^2 \theta_w + 2 \sin^4 \theta_w \right) \\ &\times \frac{\sqrt{\lambda(s)}}{s} \left(E_1 E_i + \frac{\lambda(s)}{12s} - m_{\chi_1^0} m_{\chi_i^0} \right) \end{aligned} \quad (11)$$

(note the different basis for the neutralinos in [37]) with

$$\lambda(s) = s^2 + m_{\chi_1^0}^4 + m_{\chi_i^0}^4 - 2s \left(m_{\chi_1^0}^2 + m_{\chi_i^0}^2 \right) - 2m_{\chi_1^0}^2 m_{\chi_i^0}^2. \quad (12)$$

In order to obtain an approximate expression for the resulting constraints on the parameters, we first neglect $m_{\chi_1^0}$ everywhere in (11). Using numerical values for the gauge couplings, (11) simplifies to

$$\sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0) [\text{pb}] \simeq 4.9 \times 10^4 \frac{(s - m_{\chi_i^0}^2)^2}{s(s - M_Z^2)^2} \left(1 + \frac{m_{\chi_i^0}^2}{2s} \right) (N_{13}N_{i3} - N_{14}N_{i4})^2 \quad (13)$$

with s and the masses in GeV. Next we look for approximations for the relevant neutralino mixing parameters N_{ij} . For simplification we assume $M_2, 2\kappa s \gg |\mu_{eff}|$ such that the wino- and singlino-sectors in the mass matrix \mathcal{M}_0 in Eq. (5) decouple. (The wino- and singlino-components of the LSP hardly contribute to the spin-independent cross section.) Assuming, in addition, large $\tan \beta$ such that $v_d \ll v_u$, \mathcal{M}_0 can be diagonalised analytically

with the results (we define $u = g_1 v_u / \sqrt{2} \sim 43$ GeV and write $\mu \equiv \mu_{eff}$)

$$\begin{aligned} N_{11} &\sim \frac{-1}{\sqrt{1 + \frac{u^2}{\mu^2}}}, N_{13} \sim \frac{-1}{\sqrt{1 + \frac{\mu^2}{u^2}}}, N_{14} \sim 0 \\ N_{21} &\sim -N_{31} \sim \frac{1}{\sqrt{2}\sqrt{1 + \frac{\mu^2}{u^2}}}, N_{23} \sim -N_{33} \sim \frac{-1}{\sqrt{2}\sqrt{1 + \frac{u^2}{\mu^2}}}, N_{24} \sim N_{34} \sim \frac{1}{\sqrt{2}}. \end{aligned} \quad (14)$$

Replacing these expressions into (13), using the numerical values for s and M_Z in the denominator and, notably, approximating $m_{\chi_i^0} \sim \mu$, one ends up with

$$\sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0) [\text{pb}] \simeq 8.3 \times 10^{-7} \frac{(s - \mu^2)^2 \mu^2}{u^2 + \mu^2} \left(1 + \frac{\mu^2}{2s}\right), \quad (15)$$

where s , μ and u are in GeV. Then the upper OPAL bound on $\sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0)$ of 0.05 pb becomes a lower bound on $|\mu|$ ($\equiv \mu_{eff}$),

$$|\mu_{eff}| \gtrsim 111 \text{ GeV}. \quad (16)$$

A somewhat stronger version of the OPAL bound ($\sigma_Z < 0.01$ pb) is implemented in the default version of NMSSMTools [30, 31]. We replace it by the published value of 0.05 pb [36] for our numerical analysis. From this, without any approximations, we obtain $|\mu_{eff}| \gtrsim 114$ GeV (varying somewhat with M_2 and $\tan \beta$) for small values of $m_{\chi_1^0}$ in good agreement with the previous estimation. We remark that, within the approximations used in Eqs. (14), this implies an upper bound on N_{13} of $N_{13}^2 \lesssim 0.12$.

Next, we consider constraints from the upper bound on the invisible Z decay width, to which the decay $Z \rightarrow \chi_1^0 \chi_1^0$ would contribute. From [38] we obtain $\Delta \Gamma_Z^{inv} \lesssim 2.0$ MeV (a value slightly above the one used in [10], but below the value used in [21]). The expression for the contribution to $\Delta \Gamma_Z^{inv}$ from χ_1^0 reads

$$\Delta \Gamma_Z^{inv} = \frac{M_Z^3 G_F}{12\sqrt{2}\pi} (N_{13}^2 - N_{14}^2)^2 \left(1 - \frac{4m_{\chi_1^0}^2}{M_Z^2}\right)^{3/2} \sim 0.165 \text{ GeV} (N_{13}^2 - N_{14}^2)^2, \quad (17)$$

where the last expression holds for small $m_{\chi_1^0}$. Then the upper bound on $\Delta \Gamma_Z^{inv}$ implies

$$|N_{13}^2 - N_{14}^2| < 0.11. \quad (18)$$

For large $\tan \beta$, where $N_{14}^2 \ll N_{13}^2$, this bound on N_{13}^2 is very similar to the bound obtained above from the OPAL limits. According to the numerical analysis without approximations we find that the constraints on the parameter space from $e^+e^- \rightarrow \chi_1^0 \chi_i^0$ are mostly somewhat stronger than those from $\Delta \Gamma_Z^{inv}$; from (8) and (9) it should be clear, that these constraints are relevant for upper bounds on the spin-independent and spin-dependent LSP-nucleon cross sections.

For the chargino masses we require a lower bound of 103 GeV [39], which implies lower limits on combinations of the parameters M_2 and μ_{eff} . In the neutral Higgs sector we apply the various constraints from [40]. Since the lightest CP-even Higgs boson h is mostly Standard-Model-like in our case, these constraints reduce to the well-known bound $m_h > 114$ GeV. On the other hand the constraints from B-physics, as described below, will imply charged Higgs masses above ~ 200 GeV, hence additional bounds from direct charged Higgs production are not required.

3.2 Constraints from B-physics

Relevant constraints from B-physics originate from bounds on $BR(b \rightarrow s\gamma) = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4}$ [41], $\Delta M_s = 17.77 \pm 0.12 \text{ ps}^{-1}$ [42] and $\Delta M_d = 0.507 \pm 0.005 \text{ ps}^{-1}$ [41], and the branching ratios $BR(B_s \rightarrow \mu^+\mu^-) < 5.8 \times 10^{-8}$ [43] (which was recently improved to $< 4.3 \times 10^{-8}$ at 95% C.L. [44]) and $BR(B^+ \rightarrow \tau^+\nu_\tau) < (1.67 \pm 0.39) \times 10^{-4}$ [41]. These constraints are implemented in NMSSMTools as described in [45], to which we refer for the corresponding contributions to these observables in the NMSSM.

It should be noted that charged Higgs boson exchange contributes to $BR(b \rightarrow s\gamma)$ and $BR(B^+ \rightarrow \tau^+\nu_\tau)$, hence the corresponding limits impose lower bounds on m_{H^\pm} . On the other hand, Susy diagrams also contribute to these observables which depend on parameters like M_2 , μ_{eff} , M_{squark} and A_{top} [46]. For specific choices of these parameters (notably not too large positive values of A_{top}), the charged Higgs boson contributions can be partially cancelled. This will be relevant below, since the spin-independent LSP-nucleon cross section (8) is maximal for small m_H and, as noted above, $m_H \sim m_{H^\pm}$.

At large $\tan\beta$, the observables ΔM_s , ΔM_d and $BR(B_s \rightarrow \mu^+\mu^-)$ can receive large contributions from a light CP-odd Higgs boson A_1 [45] which, in turn, plays an important rôle for the LSP annihilation cross section (10) for a small LSP mass $m_{\chi_1^0}$. Again, additional Susy contributions (box diagrams) exist, leading to a complicated combination of constraints in the parameter space. We find that, for a small LSP mass (light A_1), practically all these observables impose bounds on various corners in the parameter space.

3.3 Additional constraints

On the dark matter relic density we impose the 3σ WMAP bound [7]

$$0.091 < \Omega h^2 < 0.129, \quad (19)$$

which requires a sufficiently large LSP annihilation rate (10).

A light CP-odd Higgs boson A_1 with a mass below $\sim 9.3 \text{ GeV}$ can appear in radiative $\Upsilon \rightarrow A_1\gamma$ decays, on which CLEO [47] and BaBar [48, 49] have obtained upper bounds. These can be translated into the parameter space (couplings of A_1) of the NMSSM [50–52] and are implemented, together with constraints from possible $A_1 - \eta_b$ mixing effects [51], in NMSSMTools. We find that these constraints are so strong (imposing, essentially, strong upper bounds on the $A_1 b\bar{b}$ coupling for $m_{A_1} \lesssim 10 \text{ GeV}$) that it becomes very difficult to obtain a LSP annihilation rate compatible with (19) for $m_{\chi_1^0} \lesssim 2 \text{ GeV}$.

Finally we require that the Susy contributions to the anomalous magnetic moment of the muon (see [53, 54] for such contributions in the NMSSM) improve the disagreement between the result of the E821 experiment [55] and the Standard Model; as in the MSSM, this implies a positive value for μ_{eff} .

4 Results and conclusions

Before we turn to our results, we discuss the range of parameters used to maximise the direct detection cross sections respecting the experimental constraints above. First, for λ

we choose a small value $\lambda = 0.05$ such that its negative effect on $M_{H^\pm}^2$ as in Eq. (4) remains negligible while a non-zero doublet component of A_1 induced by the off-diagonal term in Eq. (3). A large value for $\kappa = 0.55$ makes the singlino (and the singlet-like CP-even Higgs state) heavy such that perturbing mixing effects in these sectors are avoided.

For the Susy breaking squark and slepton masses we use 1 TeV such that sleptons hardly contribute to $e^+e^- \rightarrow \chi_1^0 \chi_i^0$. A_{top} varies from 300 to 650 GeV where H^\pm -induced and Susy-induced contributions to $BR(b \rightarrow s\gamma)$ tend to cancel [45]. The Susy breaking gluino and the wino masses are chosen as $M_3 = 350$ GeV and $M_2 = 180$ GeV, respectively. (These parameters appear in the loop-induced flavour changing A_1 -quark vertices [46], which should be small in order to allow for a light A_1 consistent with the constraints from $BR(B_s \rightarrow \mu^+ \mu^-)$.)

Although Eq. (8) suggests that σ^{SI} is maximised for very large values of $\tan \beta$, the best compromise with B-physics is obtained for reasonable values of $\tan \beta \sim 35 - 44$. Likewise, Eq. (14) suggests that μ_{eff} should be as small as possible in order to maximise $N_{11}N_{13}$, but we find that the best compromise in parameter space is obtained for $\mu_{eff} \simeq 128$ GeV. Eq. (8) also suggests that σ^{SI} is maximised for m_H as small as possible. However, we recall that $m_H \sim M_A \sim m_{H^\pm}$ and that m_{H^\pm} is bounded from below by several B-physics processes. We choose M_A as an input parameter of the NMSSM (instead of A_λ) and find the largest values of σ^{SI} for $M_A \sim 260 - 315$ GeV, implying $m_H \sim 205 - 260$ GeV and $m_{H^\pm} \sim 225 - 280$ GeV where the larger values correspond to lower masses of $m_{\chi_1^0}$ below.

The Susy breaking bino mass term M_1 determines $m_{\chi_1^0}$, with $M_1 \sim 23.5$ GeV for $m_{\chi_1^0} \sim 20$ GeV and $M_1 \sim 3.0$ GeV for $m_{\chi_1^0} \sim 2.0$ GeV. Finally A_κ is chosen in the range $A_\kappa \sim -14 \dots -4$ GeV, which determines m_{A_1} such that χ_1^0 pair annihilation via A_1 exchange in the s-channel gives the correct relic density in agreement with the WMAP bound in Eq. (19). Due to the relatively large couplings involved in the χ_1^0 pair annihilation process, A_1 must actually be off-shell and hence m_{A_1} substantially larger than $2m_{\chi_1^0}$; otherwise the relic density is too small. In Fig. 1 we show the result for m_{A_1} as function of $m_{\chi_1^0}$.

For $m_{A_1} \lesssim 40$ GeV ($m_{\chi_1^0} \lesssim 5$ GeV) the constraints from $BR(B_s \rightarrow \mu^+ \mu^-)$ (where A_1 appears in the s-channel) become particularly strong and require a somewhat smaller doublet component of A_1 . Denoting its doublet component by $\sin \theta_A$, we have $\sin \theta_A \sim 0.8$ for $m_{A_1} \gtrsim 50$ GeV, but $\sin \theta_A \sim 0.45$ for $m_{A_1} \sim 10$ GeV. We note that for $m_{A_1} \lesssim 40$ GeV the value of A_κ has to be chosen within a precision less than 1% such that the relic density of χ_1^0 is below the WMAP bound (possibly smaller), but $m_{A_1}^2 > 0$; hence this region in the parameter space requires considerable fine tuning. For $m_{A_1} < 10$ GeV ($m_{\chi_1^0} \lesssim 2$ GeV) the constraints from CLEO and BaBar become so strong that $\sin \theta_A$ must be much smaller requiring an even stronger fine tuning of parameters, therefore we will not consider this range of parameters subsequently.

The components of χ_1^0 (the coefficients N_{1i} , see Eq. (6)) hardly change in the range $m_{\chi_1^0} = 2 - 20$ GeV considered here, once we maximise the product $N_{11}N_{13}$ in order to maximise σ^{SI} . We have

$$\begin{aligned} N_{11} &\sim -0.94, \quad N_{12} \sim 0.01 \dots 0.03, \quad N_{13} \sim -0.32 \dots -0.34, \\ N_{14} &\sim 0.013 \dots 0.06, \quad N_{15} \sim 0.001. \end{aligned} \tag{20}$$

The masses of the mostly higgsino-like states χ_2^0 and χ_3^0 are ~ 105 and ~ 145 GeV, respec-

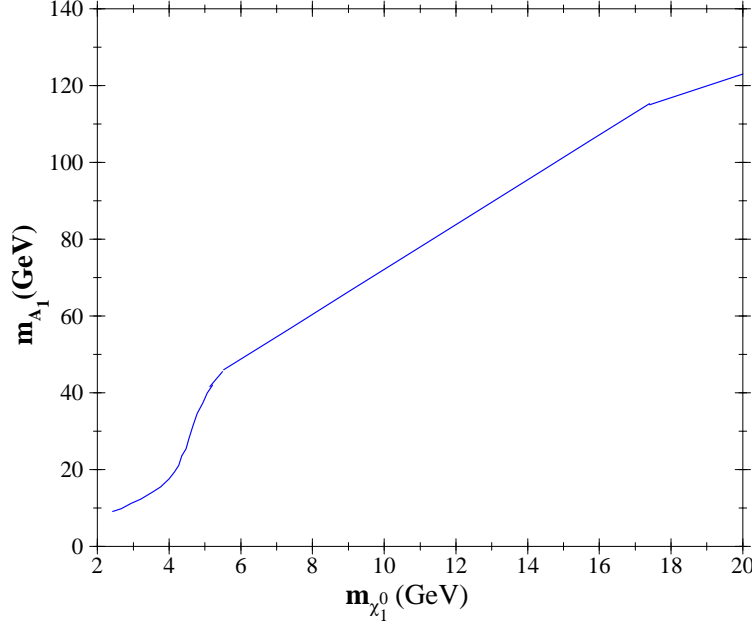


Figure 1: m_{A_1} as function of $m_{\chi_1^0}$ such that the relic density of χ_1^0 is in agreement with the WMAP bound Eq. (19).

tively, and hence as stated before, the limits on $\sigma_Z(e^+e^- \rightarrow \chi_1^0 \chi_i^0)$ are relevant.

The scattering rates of χ_1^0 depend somewhat on astrophysical parameters as the escape velocity v_{max} and the dark matter density ρ_0 near the sun and, more importantly, on nuclear form factors (quark matrix elements) as the pion-nucleon sigma term $\sigma_{\pi N}$ and the size of $SU(3)$ symmetry breaking parametrized by σ_0 . (The difference $\sigma_0 - \sigma_{\pi N}$ is proportional to the strange quark matrix element.) For the astrophysical parameters we use the default values of micrOMEGAs $v_{max} = 544$ km/s and $\rho_0 = 0.3$ GeV/cm³ [34]. The default values in micrOMEGAs for $\sigma_{\pi N}$ and σ_0 are $\sigma_{\pi N} = 55$ MeV and $\sigma_0 = 35$ MeV.

The corresponding results for the upper bound on the spin-independent cross section of χ_1^0 off protons σ_p^{SI} in the NMSSM are shown in Fig. 2 as a function of $m_{\chi_1^0}$ as a full red line. (The spin-independent cross section off neutrons is nearly the same.) In order to indicate the variation of this upper bound with $\sigma_{\pi N}$ and σ_0 , we show a red dashed line as the upper bound on σ_p^{SI} for $\sigma_{\pi N} = 73$ MeV and $\sigma_0 = 30$ MeV, which would correspond to a larger strange quark matrix element and hence an increase of σ^{SI} by a factor ~ 3.3 .

Also shown in Fig. 2 are the regions compatible with the excesses of events reported by DAMA [1] (without channeling (dark blue) and with channeling (light blue)), CoGeNT [2] (light green) and a fit to the two events observed by CDMS-II [56] (denoted as CDMS-09 fit surrounded in dashed green; these events are also compatible with background). Exclusion limits are shown from Xenon10 [4] (violet), Xenon100 [5] (black) and CDMS-II [3, 6] (magenta, assuming that the two observed events originate from background).

Fig. 2 is our main result, which leads to the following conclusions:

- It seems difficult to explain the excesses of events reported by DAMA and CoGeNT within the general NMSSM (without unification constraints on M_1). Hence, as stated

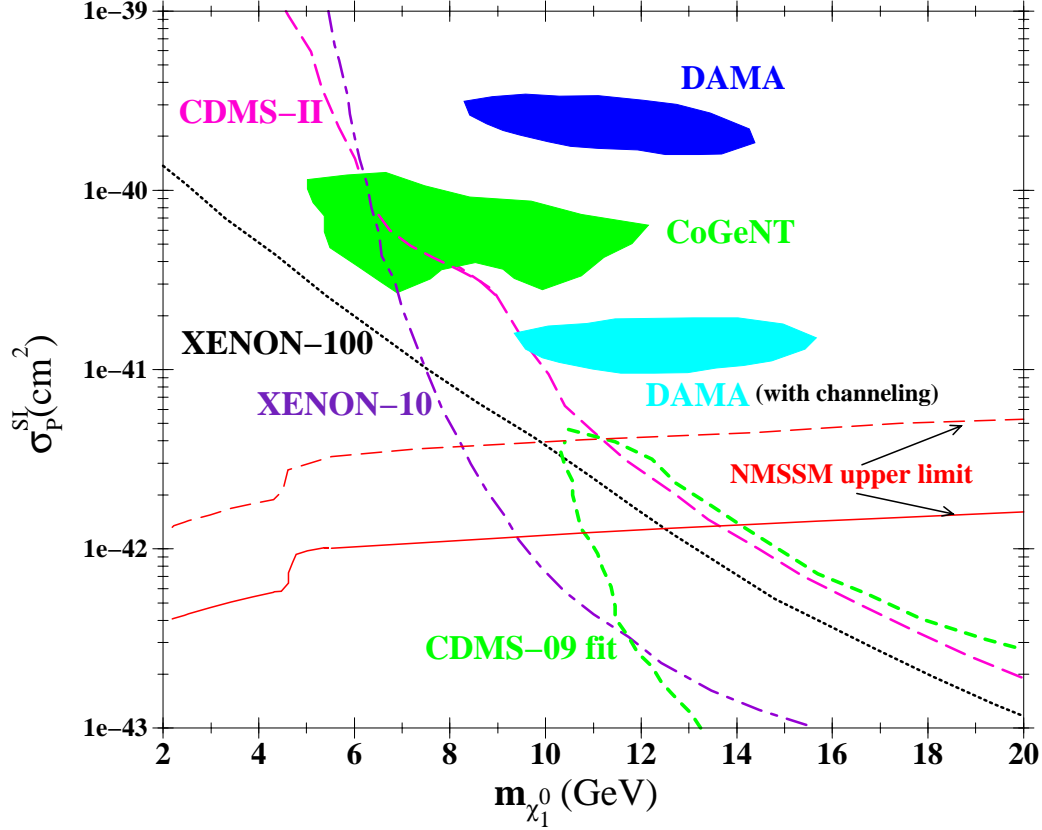


Figure 2: Upper bounds on the spin-independent cross section σ_p^{SI} in the NMSSM for default values of the strange quark content of nucleons as a full red line, and an enhanced strange quark content of nucleons as a dashed red line. Also shown are regions compatible with DAMA, CoGeNT and CDMS-II, and limits from Xenon10, Xenon100 and CDMS-II as explained in the text.

in [21], significant modifications of parameters like a larger local dark matter density ρ_0 would be required to this end. On the other hand, the two events observed by CDMS-II (within the contour denoted as CDMS-09 fit) could be explained in the NMSSM.

- Actual limits of Xenon10, Xenon100 and CDMS-II on spin-independent cross sections of WIMPS in the 2 – 20 GeV mass range test regions of the parameter space of the NMSSM.

For completeness we have also considered the spin-dependent cross section σ^{SD} in the NMSSM, which is maximal for $\tan \beta \gtrsim 20$ (such that $N_{14}^2 \ll N_{13}^2$ in Eq. (9)), large values of M_A (since m_H is irrelevant here), and $\mu_{eff} \sim 121 - 129$ GeV. In Fig. 3 we show the maximum of the spin-dependent cross section off protons σ_p^{SD} for the same range of $m_{\chi_1^0} = 2 - 20$ GeV. Note that σ^{SD} originates from Z -exchange, hence the spin-dependent cross section off neutrons σ_n^{SD} is given by $\sigma_n^{SD} \simeq 0.78 \times \sigma_p^{SD}$. The actual experimental upper

limits on σ^{SD} are one to two orders of magnitude larger [57] than the upper bounds in the NMSSM and not shown in Fig. 3.

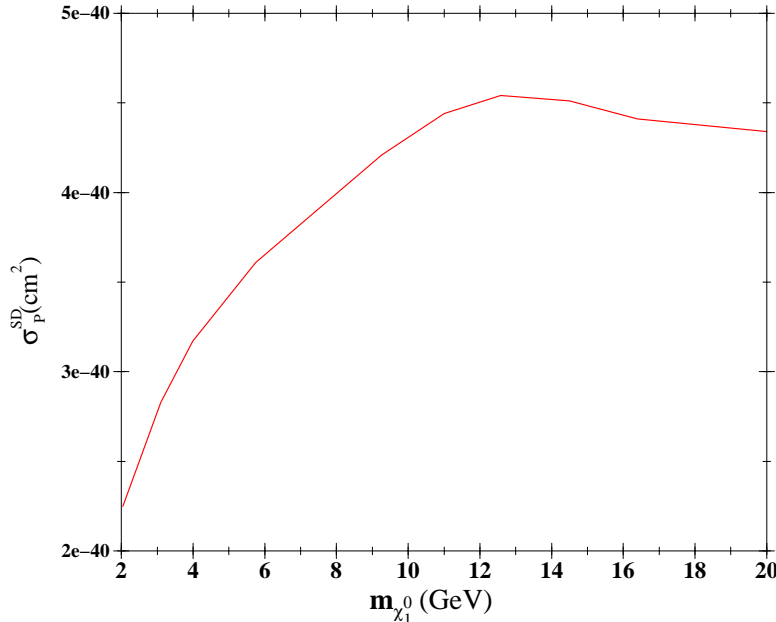


Figure 3: Upper bounds on the spin-dependent cross section σ_p^{SD} in the NMSSM.

To conclude, we have performed a detailed analysis of the parameter space of the NMSSM for general values of M_1 , which allows for WIMP masses in the 2 - 20 GeV range. In contrast to the MSSM, light bino-like WIMPs can have a relic density compatible with WMAP constraints due to a light NMSSM-specific CP-odd Higgs state which can be exchanged in the s-channel. Due to reported excesses of events compatible with WIMP masses below 20 GeV, this region is of particular interest.

We have studied in detail the constraints on this region of the parameter space of the NMSSM from LEP and B-physics, and the regions of parameter space which give rise to maximal direct detection cross sections while not contradicting experimental limits. The resulting upper bounds on $\sigma^{SI} \lesssim 10^{-42} \text{ cm}^2 = 10^{-6} \text{ pb}$ make it difficult to explain the excesses of events reported by DAMA and CoGeNT within the NMSSM for small values of M_1 . On the other hand, the two events observed by CDMS-II could be explained in the NMSSM.

Notably the Xenon10 limits [4] on σ^{SI} for WIMP masses below 20 GeV start to test corresponding regions of the NMSSM parameter space. Future results from Xenon100 could confirm the presence of a light WIMP compatible with the NMSSM, or impose further constraints on its parameter space.

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